

R –Parity Violation in a Warped GUT Scale Randall-Sundrum FrameworkB.C. Allanach,^{1,*} A. M. Iyer,^{2,†} and K. Sridhar^{2,‡}¹*DAMTP, CMS, Wilberforce Road, University of Cambridge,
Cambridge, CB3 0WA, United Kingdom*²*Department of Theoretical Physics, Tata Institute of Fundamental Research,
Homi Bhabha Road, Colaba, Mumbai 400 005, India***Abstract**

We consider a modified Randall-Sundrum (RS) framework between the Planck scale and the GUT scale. In this scenario, RS works as a theory of flavour and not as a solution to the hierarchy problem. The latter is resolved by supersymmetrising the bulk, so that the minimal supersymmetric standard model being the effective 4-dimensional theory. Matter fields are localised in the bulk in order to fit fermion-mass and mixing-data. If R -parity violating (\mathcal{R}_p) terms are allowed in the superpotential, their orders of magnitude throughout flavour space are then predicted, resulting in rich flavour textures. If the \mathcal{R}_p contributions to neutrino masses are somewhat suppressed, then lepton-number violating models exist which explain the neutrino oscillation data while not being in contradiction with current experimental bounds. Another promising model is one where baryon number is violated and Dirac neutrino masses result solely from fermion localisation. We sketch the likely discovery signatures of the baryon-number and the lepton-number violating cases.

*Electronic address: b.c.allanach@damtp.cam.ac.uk

†Electronic address: abhishek@theory.tifr.res.in

‡Electronic address: sridhar@theory.tifr.res.in

I. INTRODUCTION

The discovery of the Higgs boson by the CMS [1] and ATLAS [2, 3] collaborations at the Large Hadron Collider has validated the status of the Standard Model (SM) as the correct theory of nature at the electroweak scale. The existence of a fundamental scalar in the SM raises questions regarding the stability of the Higgs mass in the face of radiative corrections. Supersymmetry emerges as one of the most exciting prospects to address this problem due to its renormalizable nature and its consistency with electroweak precision data. The model, however, introduces a number of additional parameters which necessitates the study of its phenomenological implications using simplified models.

However, supersymmetry in its minimal form *viz* the minimal supersymmetric standard model (MSSM) does not give an explanation of the disparate couplings of the Higgs boson to different generations of fermions. This is also referred to as the fermion mass problem. In supersymmetry this problem can be addressed by considering strong wave function renormalization of the matter fields [4, 5]. This is due to renormalization group (RG) running from some fundamental renormalization scale (where the Yukawa as well as the soft mass parameters are anarchic) to a low scale where they develop a hierarchical structure due to renormalization effects. The different running can be accounted by the different anomalous dimensions of each matter field coupled to some conformal sector. After canonically normalizing the kinetic terms, the superpotential terms are given by

$$\begin{aligned} \mathcal{W} = & \epsilon^{q_i+h_u+u_j} \left(\hat{Y}_U^{ij} + \hat{A}_U^{ij} X \right) Q_i H_u U_j + \epsilon^{q_i+h_u+d_j} \left(\hat{Y}_D^{ij} + \hat{A}_D^{ij} X \right) Q_i H_d D_j \\ & + \epsilon^{l_i+h_d+e_j} \left(\hat{Y}_e^{ij} + \hat{A}_e^{ij} X \right) L_i H_d E_j \end{aligned} \quad (1)$$

while the Kähler terms are given by

$$\mathcal{K} = \sum_{F=Q,U,D,L,E} F^\dagger F + \hat{C}_{ij} \epsilon^{f_i+f_j} X^\dagger X F^\dagger F \quad (2)$$

where i and j are generation indices and quantities with hatted quantities denoting $\mathcal{O}(1)$ parameters. X is the SUSY breaking spurion parametrized as $X = \theta^2 F$. The expansion parameter $\epsilon \sim 0.02$ while the ‘charges’ $q_i, h_{u,d}$ can be considered to be anomalous dimensions of the matter field coupling to a strong sector. Alternatively, they can be considered to be charges of the field under an extended gauge group $U(1)_{FN}$ [6]. The fermion mass matrix is then given as

$$m_f \sim \epsilon^{q_i+h_u+u_j} \frac{v}{\sqrt{2}} \quad (3)$$

where $v \sim 246$ GeV is the vacuum expectation value (VEV) of the Higgs field. The charges q_i are determined with the requirement of reproducing the correct pattern of fermion mass and

mixing angles. Soft supersymmetry breaking terms are generated when F terms attains a VEV giving rise to the gravitino mass $m_{3/2} = \frac{\langle F \rangle}{M_{Pl}}$. The mechanism which fixes the fermion masses and mixing angles will also determine the soft supersymmetry-breaking mass parameters, for example the squark mass-squared terms:

$$\tilde{m}_{ij}^2 \sim \mathcal{O}(\epsilon^{q_i+q_j} m_{3/2}^2), \quad (4)$$

and similarly for the other family-dependent soft supersymmetry-breaking terms. This lends a certain level of predictivity to the orders of magnitude of soft breaking terms.

The terms in Eqs. 1,2 conserve R -parity [7, 8], which is defined as

$$R = (-1)^{3B+L+2s} \quad (5)$$

where s is spin of a particle and $B(L)$ is its corresponding baryon (lepton) number (alternatively, the same terms conserve matter parity [9–11]). While R parity conservation has many useful features, predicting the stability of dark matter and a stable proton, there is no *a priori* reason for it to be a symmetry of the lagrangian¹. Thus in general, the super-potential terms in Eq.(1) can also be extended to include terms which violate baryon and lepton-number, and are referred to as R -parity violating (\mathcal{R}_p) terms². The most general \mathcal{R}_p terms are given by

$$\begin{aligned} \mathcal{W}_{\mathcal{R}_p}^{\Delta L=1} &= \frac{\epsilon^{l_i+l_j+e_k}}{2} \lambda_{ijk} L_i L_j E_k + \frac{\epsilon^{l_i+q_j+d_k}}{2} \lambda'_{ijk} L_i Q_j D_k + \epsilon^{l_i+h_u} \mu'_i L_i H_u, \\ \mathcal{W}_{\mathcal{R}_p}^{\Delta B=1} &= \frac{\epsilon^{u_i+d_j+d_k}}{2} \lambda''_{ijk} U_i D_j D_k, \end{aligned} \quad (6)$$

where we have omitted the gauge indices. In fact, it is easy to find [17] anomaly-free symmetries which ban a some subset of Eq. 6: baryon triality sets $\mathcal{W}_{\mathcal{R}_p}^{\Delta B=1} = 0$ whereas lepton parity bans $\mathcal{W}_{\mathcal{R}_p}^{\Delta L=1}$. These symmetries could be separately imposed upon the theory, and we shall take advantage of this fact later.

Whenever R -parity violation is introduced, one wonders where the apparent relic density of dark matter might come from, given that it appears to be stable on cosmological scales, and any MSSM fields will decay much too quickly. One obvious answer is that massive hidden sector matter, might provide dark matter. Unfortunately, this would result in *no* direct or indirect signals for dark matter detection. Another possibility [19, 20] is that the lightest supersymmetric particle is the gravitino, which has Planck suppressed couplings anyway. With additional smallish \mathcal{R}_p violating

¹ Scenarios in which R parity originated as a discrete remnant of some extended gauge symmetries were considered in [12–17].

² For a detailed review on (\mathcal{R}_p) supersymmetry see Ref. [18].

couplings, it is possible that its lifetime is much longer than the age of the universe, resulting in a good dark matter candidate. We leave this aspect of the model building to a future paper.

The anomalous dimensions of the matter fields can also be considered dual to the parameter which controls the localization of the field in an extra-dimensional scenario with strong warping [21–23]. In this paper we consider the effects of introducing all such terms in a supersymmetric model on a gravitational background with strong warping also referred at as Randall-Sundrum (RS) model [24]. In Section II we briefly introduce the model and set it up to understand the phenomenology. We review the technique used to determine the RS-model parameters which fit the fermion-mass and mixing-data at the high scale. The mathematical expressions used to determine the soft- and \hat{R}_p -parameters are presented. In section III we discuss the implications of introducing \hat{R}_p couplings on various low-energy processes. We find that if baryon-number and lepton-number violating terms are simultaneously allowed, consistency with constraints from proton decay require a slightly fine tuned choice of 10^{-4} in some undetermined parameters usually expected to be of $\mathcal{O}(1)$. We then proceed to discuss simplified cases where either baryon- or lepton-number is conserved separately. Scenarios with lepton-number violation present solutions where the neutrinos can be Dirac-like even in the presence lepton number violating terms. In each case, we briefly comment on the LHC phenomenology, before presenting our conclusions.

II. GUT-SCALE RANDALL-SUNDRUM MODEL

We consider the following modified version of the original setup referred to as ‘GUT-scale RS’ [5, 25–29]. Like the original RS model, it is a model of single extra-dimension compactified on a S_1/Z_2 orbifold. The line-element is given as

$$ds^2 = e^{-2\sigma(y)} \eta_{\mu\nu} dx^\mu dx^\nu + dy^2 \quad (7)$$

where $\sigma(y) = k|y|$ with k denoting the reduced Planck scale and $R \sim 1/k$ being the size of the extra spatial dimension y . There are two opposite tension branes at each of the orbifold fixed points, $y = 0$ and $y = \pi R$. Assuming the scale of physics at the $y = 0$ brane to be M_{Pl} , the effective scale induced at the brane at $y = \pi R$ is given by

$$M_{IR} = e^{-\sigma(\pi)} M_{Pl} = \epsilon M_{Pl} \sim M_{GUT} \quad (8)$$

Thus, in comparison to the original proposal in Ref. [24], the warp factor in this case is much larger and hence *ab initio* the model is no longer a solution to the hierarchy problem. Hence, supersym-

metry is introduced into the bulk. With the GUT-scale Kaluza Klein (KK) modes decoupled from the theory, the spectrum of the effective 4D theory is that of MSSM.

We assume the two Higgs doublets to be localized on the infra-red (IR) brane (i.e. on the GUT brane) while the matter and gauge multiplets are in the bulk. The expressions for the fermion mass matrices are

$$\begin{aligned}(m_u)_{ij} &= v_u \hat{Y}_{u_{ij}} f(c_{Q_i}) f(c_{u_j}) \\ (m_d)_{ij} &= v_d \hat{Y}_{d_{ij}} f(c_{Q_i}) f(c_{d_j}) \\ (m_e)_{ij} &= v_d \hat{Y}_{e_{ij}} f(c_{L_i}) f(c_{e_j})\end{aligned}\tag{9}$$

where i and j are generation indices, $v_u = \frac{v}{2} \sin \beta$ and $v_d = \frac{v}{2} \sin \beta$ are the VEVs of the up-type and down-type Higgs' of the MSSM, respectively and c_{Z_i} are the dimensionless bulk mass parameters of the matter multiplets $Z \in \{Q, u, d, L, e\}$. The corresponding zero-mode wave-function f is defined to be [21, 30]

$$f(c) = \sqrt{\frac{1-2c}{e^{(1-2c)\pi k R} - 1}} e^{(0.5-c)k R \pi}.\tag{10}$$

Using Eq. 9 and choosing $c_{Z_i} \sim \mathcal{O}(1)$ and $\hat{Y}_{U,D,E}^{ij} \sim \mathcal{O}(1)$, one can explain the observed hierarchy in the fermion masses and mixings [21, 29–32].

A SUSY-breaking spurion $X = \theta^2 F$ is introduced on the GUT brane and IR brane-localized contact interactions are introduced between the SM fields and the SUSY breaking spurion X ³. The soft SUSY breaking terms are then generated when the F -term attains a VEV and are given by

$$\begin{aligned}m_{H_{u,d}}^2 &= \hat{h}_{u,d} m_{3/2}^2 \\ (m_Z^2)_{ij} &= m_{3/2}^2 \hat{\beta}_{\tilde{Z}_{ij}} f(c_{Z_i}) f(c_{Z_j}) \\ A_U^{ij} &= m_{3/2} \hat{A}_U^{ij} f(c_{Q_i}) f(c_{u_j}), \\ A_D^{ij} &= m_{3/2} \hat{A}_D^{ij} f(c_{Q_i}) f(c_{d_j}), \\ A_E^{ij} &= m_{3/2} \hat{A}_E^{ij} f(c_{L_i}) f(c_{e_j}), \\ m_\alpha &= \hat{g}_\alpha m_{3/2},\end{aligned}\tag{11}$$

where quantities denoted with a hat are fundamental dimensionless parameters, which we assume are $\mathcal{O}(1)$. Here, $m_{H_{u,d}}^2$ are the up- and down- Higgs mass squared soft SUSY-breaking parameters,

³ For supersymmetric version of regular RS models see for example [21, 25, 33, 34]

$(m_{\tilde{Z}}^2)_{ij}$ the soft SUSY breaking mass squared matrix for sfermion \tilde{Z} , $A_{(U,D,E)}^{ij}$ the matrix of trilinear soft SUSY-breaking interactions for the up-quark, down-quark and charged leptons and m_i the i^{th} gaugino mass (where $\alpha \in \{3, 2, 1\}$ denotes the MSSM gauge group $SU(3)$, $SU(2)_L$, $U(1)_Y$), respectively. Thus, the soft masses are flavourful, i.e. there is flavour mixing between different generations of sfermions. However, as Eq. 11 shows, the fermion masses and soft masses are determined by the same set of c_i parameters. Since both the Higgs superfields and the SUSY breaking spurion are localized on the same brane, superfields with small fermion masses have small scalar masses. As a result, masses of the lighter generations are very small at the GUT scale and the masses of the lighter generation sfermions at the weak scale are almost entirely governed by the gaugino masses due to renormalisation effects. This is the reason why the first two generations are almost degenerate.

R -parity violating interactions are introduced on the GUT brane and so they are considered to be generated at the GUT scale⁴. Like the soft parameters, the effective four-dimensional (4D) \mathcal{R}_p parameters can also be expressed in terms of the bulk wavefunction of the fields. The effective 4D \mathcal{R}_p violating superpotential in a warped background is written

$$\begin{aligned}\mathcal{W}_{\mathcal{R}_p}^{\Delta L=1} &= \int dy e^{-3ky} \delta(y - \pi R) \left(\lambda_{ijk}^{(5)} L_i L_j E_k + \frac{1}{2} \lambda'_{ijk}{}^{(5)} L_i Q_j D_k + \mu_i^{(5)} L_i H_u \right) \\ \mathcal{W}_{\mathcal{R}_p}^{\Delta B=1} &= \int dy e^{-3ky} \delta(y - \pi R) \lambda''_{ijk}{}^{(5)} U_i D_j D_k.\end{aligned}\quad (12)$$

$\lambda_{ijk}^{(5)}, \lambda'_{ijk}{}^{(5)}, \lambda''_{ijk}{}^{(5)}$ are 5D \mathcal{R}_p couplings which have mass dimension -1 . Performing a KK decomposition of the fields and retaining only the zero modes⁵, the effective 4D \mathcal{R}_p couplings are written as

$$\begin{aligned}\lambda_{ijk} &= \hat{\lambda}_{ijk} f(c_{L_i}) f(c_{L_j}) f(c_{E_k}) \\ \lambda'_{ijk} &= \hat{\lambda}'_{ijk} f(c_{L_i}) f(c_{Q_j}) f(c_{D_k}) \\ \mu_i &= \mu \hat{\mu}_i f(c_{L_i}) e^{-kR\pi}\end{aligned}\quad (13)$$

for the $\Delta L = 1$ terms. μ is of order the electroweak scale and is chosen here to be 100 GeV. $\hat{\lambda}_{ijk}, \hat{\lambda}'_{ijk}, \hat{\mu}_i$ are dimensionless $\mathcal{O}(1)$ couplings, where as $\hat{\lambda} \equiv k\lambda_{ijk}^{(5)}$, $\hat{\lambda}' \equiv k\lambda'_{ijk}{}^{(5)}$ and $\hat{\mu} \equiv k\mu^{(5)}$. The $\Delta B = 1$ \mathcal{R}_p couplings are

$$\lambda''_{ijk} = \hat{\lambda}''_{ijk} f(c_{U_i}) f(c_{D_j}) f(c_{D_k}) \quad (14)$$

⁴ An equivalent description would correspond to the Higgs doublets and the \mathcal{R}_p violating terms localized on the ultra-violet (UV) brane.

⁵ Higher KK modes in this model have mass $\sim M_{GUT}$ and are decoupled.

with $\hat{\lambda}_{ijk}'' = k\lambda_{ijk}''^{(5)}$. The supersymmetric parameters in Eqs. 11,13 and 14 are determined using the set of same set of c_i parameters that fit the fermion masses and mixing at the GUT scale using Eqs. 9 and 10. This gives an order-of-magnitude level of predictability for this framework, as these high-scale parameters can be subsequently evolved to generate a characteristic spectrum at the low scale. The set of $\mathcal{O}(1)$ parameters (which includes the c_i parameters as well as the $\mathcal{O}(1)$ Yukawa parameters $\hat{Y}_{U,D,E}^{ij}$) is determined by performing a χ^2 fit of their GUT-scale values to the data [29]. The χ^2 function is defined as

$$\chi^2 = \sum_i \frac{\left(\mathcal{O}_i^{\text{theory}}(\{c_j\}, \{\hat{Y}_U^{ij}\}, \{\hat{Y}_D^{ij}\}, \{\hat{Y}_E^{ij}\}) - \mathcal{O}_i^{\text{expt}} \right)^2}{\sigma_i^2}, \quad (15)$$

where $\mathcal{O}_i^{\text{theory}}$ denotes the theoretical prediction for observable \mathcal{O}_i , $\mathcal{O}_i^{\text{expt}}$ denotes the empirical central value and the experimental uncertainty is written σ_i . $i \in \{m_u, m_d, m_c, m_s, m_t, m_b, |V_{CKM}|_{ij}^{\text{ind}}\}$ constitute the hadronic observables whereas $i \in \{m_e, m_\mu, m_\tau, |V_{PMNS}|_{ij}^{\text{ind}}\}$ constitute the leptonic observables. Both are fit independently ('ind' indicates that the absolute values of a selection of *independent* entries - the off-diagonal entries - of the CKM and PMNS matrices are fit, respectively). We refer the interested reader to Ref. [29] for further details.

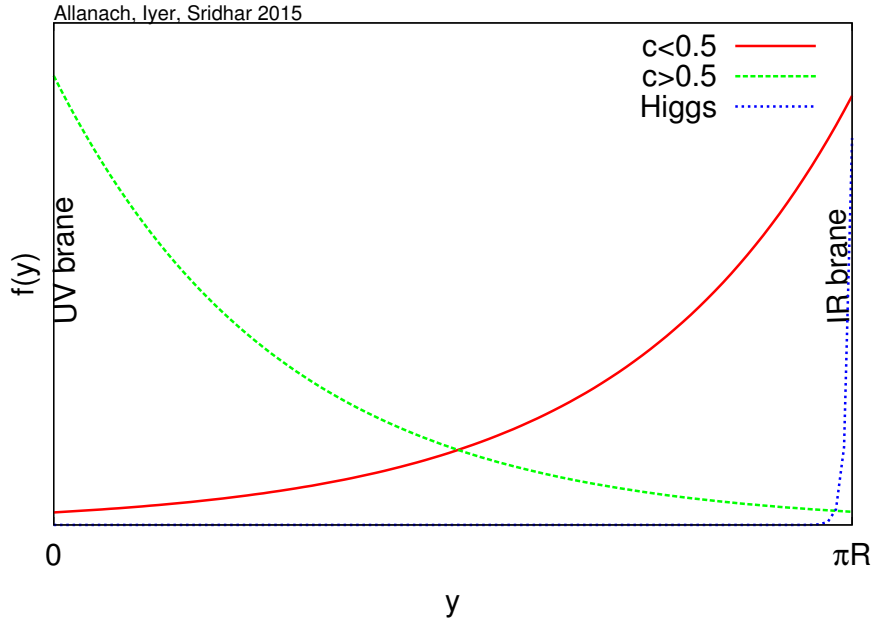


FIG. 1: Localization of fermion profiles in the bulk depends upon the c parameter. The Higgs is assumed to be strongly localized on the IR brane, as shown [35].

Because of the small value of the warp factor $\epsilon \sim 0.02$, the c_i parameters for the lighter generations (eg. the electron or the up and down quark) are close to 2.5 while for the third generation

of superfields containing a top, they are close to -1. $c_i < 0.5$ reflects a localization more on the IR brane (where the Higgs doublets are localized), while $c_i > 0.5$ localizes the superfield closer to the UV brane, as heuristically depicted in Fig. 1. The c_i parameters for all charged matter superfields except for Q^3 and t_R are scanned in the range $[0, 3.5]$, while c_{Q^3} is scanned in $[0, 1.5]$ and c_{t_R} is restricted to $[-2.5, 0.5]$. This different choice for c_{Q^3}, c_{t_R} is to facilitate a good fit to the top quark mass. We remind the reader that fits to fermion mass and mixing data are done *independently* for the quark sector and the leptonic sector, since to a good approximation (i.e. at tree-level), the two sectors are decoupled.

The fits in the leptonic sector includes fitting the neutrino data by means of introduction of three parameters c_{N_i} corresponding to three right handed neutrinos. To account for small neutrino masses at the sub eV level, c_{N_i} are scanned in the range $[5.5, 7]$ in order to imbue Dirac neutrino masses via Eq. 9. The presence of lepton-number violating operators gives rise to additional Majorana contributions to the neutrino masses. By focussing on regions of the parameter space where these contributions are suppressed, we will find that the dominant contribution to the neutrino mass is from Eq. 9 and hence are primarily of Dirac type.

The fits are performed for two separate values of $\tan \beta = 5, 10$. Smaller $\tan \beta$ facilitates a localization of the light down sector fields closer to the UV brane owing to a larger value of $\cos \beta$ in the mass matrix in Eq. 9. As we shall explain later, this helps in generating a smaller value for the \mathbb{R}_p couplings, enabling them to satisfy experimental constraints more easily, which are typically upper bounds. The $\mathcal{O}(1)$ model parameters are determined by minimising the χ^2 function in Eq. 15. The minimisation is performed by MINUIT [36] which looks for a minimum around a guess value of c parameters and $\mathcal{O}(1)$ Yukawa parameters. The guess values are randomly generated in the ranges given above. This is repeated for 10^5 choices of guess values each constituting a separate minimization. MINUIT has trouble searching our parameter space, and finds many distinct local minima, depending upon which random guess we start with. We view this as a sampling of the ‘good-fit’ parameter space, and all points which satisfy $\chi^2 < 10$ are accepted as being a reasonable ‘fit’. We remind the reader that this is not a fit to data in the usual statistical sense: rather it is a fit to the orders of magnitude of the masses and mixings. In addition to the c_i parameters, the R -parity conserving hatted $\mathcal{O}(1)$ Yukawa parameters are all allowed to vary between 0.1 and 10, whereas the hatted \mathbb{R}_p violating parameters are set to 1 and are not varied. With the c_i ’s fixed in this manner, for each sampling, we predict the orders of magnitude of the \mathbb{R}_p parameters.

III. \mathcal{R}_p PARAMETERS

We now focus on the distribution of the various \mathcal{R}_p couplings which are determined from the fermion mass fits. As given in Eq. 6, \mathcal{R}_p terms include both baryon-number and lepton-number violating interactions. The lepton-number violating interactions include the trilinear couplings $(\lambda_{ijk}, \lambda'_{ijk})$ and the bilinear operators μ_i . λ_{ijk} is anti-symmetric in $i \leftrightarrow j$ because of the $SU(2)_L$ structure, as is λ''_{kij} because of the implicit $SU(3)$ structure.

On account of the introduction of the \mathcal{R}_p operators on the same brane as the Higgs superfields, their magnitude can be roughly understood from the generation indices in these couplings. For instance, consider λ_{111} which is a product of the zero mode profiles of some first-generation fermions and λ_{333} is the corresponding product of third-generation fermions. Since, as Fig. 1 illustrates, the lighter fermion generations have a tendency to be localized away from the Higgs ($c_i > 0.5$), the corresponding value of the profile on the IR brane is small. The third generation is relatively heavy and has a value of c_i smaller than those for the lighter generations. As a result the corresponding value of the profile on the IR brane is relatively larger. This results in a larger value for λ_{333} as compared to λ_{111} . Similarly, $\mu_3 \geq \mu_2 \geq \mu_1$. In evaluating the \mathcal{R}_p parameters, the $\mathcal{O}(1)$ parameters $\hat{\lambda}_{ijk}, \hat{\lambda}'_{ijk}, \hat{\lambda}''_{ijk}$ and $\hat{\mu}_i$ were all chosen to be 1 (unless they are set to zero by requiring baryon or lepton number conservation). We therefore should bear in mind that we provide order of magnitude predictions, which will be multiplied by some order one parameter.

The predictions thus obtained are then filtered against 2 σ upper bounds on \mathcal{R}_p violating parameters, for instance from the non-observation of $\mu \rightarrow e\gamma$ [37], leptonic decay of long lived neutral Kaon [38], bounds from $n - \bar{n}$ oscillations and double neutron decay [39] or constraints from the electroweak precision tests [40–42] *etc.* A complete list of constraints on the various \mathcal{R}_p parameters that we use is given in Tables 6.1 to 6.5 of Ref. [18], although in the first instance we do not apply bounds from nucleon decay, upon which more later. The constraints do depend upon the supersymmetric spectrum, for example the branching ratio of $B \rightarrow \tau\nu$

$$\lambda'_{333} < 0.32 \left(\frac{m_{\tilde{b}_R}}{100 \text{ GeV}} \right) \quad (16)$$

depends upon the right-handed sbottom mass $m_{\tilde{b}_R}$. We shall provide predictions for viable ranges of \mathcal{R}_p violating parameters for soft masses $\tilde{m} \gtrsim 300$ GeV. If any one of the 2σ bounds is violated in the case with all hatted \mathcal{R}_p violating parameters fixed to one and 300 GeV sparticles, the point is discarded. After this filtering, we obtain 2203 good-fit parameter points to the quark mass and mixing data and 848 to the lepton mass and mixing data. We combine each set of good-fit quark

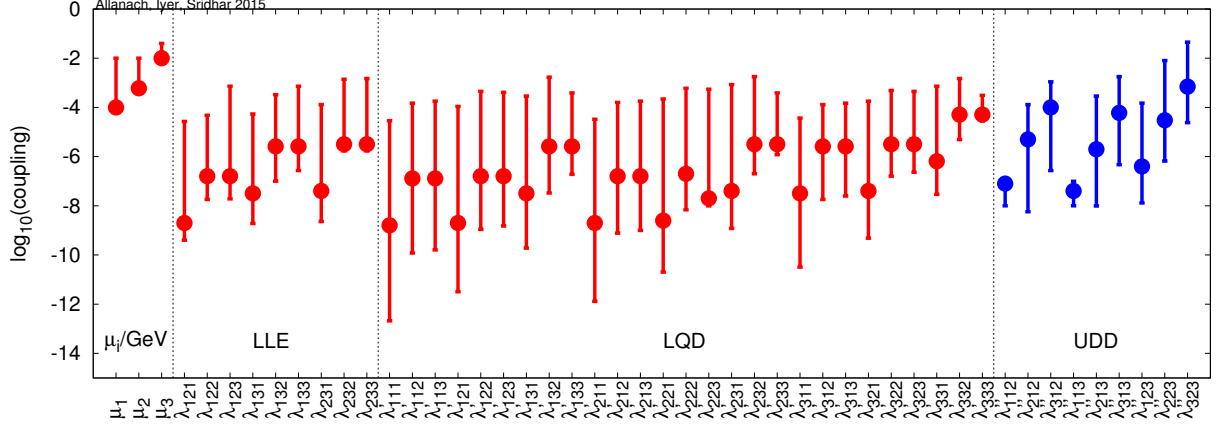


FIG. 2: Pattern of \mathcal{R}_p couplings. The vertical bars give the range of couplings that result from good fits to fermion masses and mixings for $\tan\beta = 5$, and that respect experimental bounds on \mathcal{R}_p couplings (if either only the baryon number violating or lepton-number violating couplings are allowed). The points correspond to a pattern from one particular fit (see sections III B, III C for details).

parameters with each set of good-fit lepton parameters (since they are approximately independent, as explained above) in order to determine the possible ranges of the various \mathcal{R}_p violating couplings.

Fig. 2 gives the ranges of the \mathcal{R}_p couplings predicted from the good-fit scanned points, and constitutes the main result of the present paper. Dimensionless \mathcal{R}_p couplings that are larger than around 10^{-6} result in prompt decays of the lightest supersymmetric particle at colliders, whereas if all couplings are smaller than 10^{-6} , displaced couplings result. We note that the smallest couplings are always predicted to be larger than this lower limit and so \mathcal{R}_p decays are prompt.

We note from Fig. 2 that the λ'_{ijk} couplings have a possibility to be smaller than the λ_{ijk} and the λ''_{ijk} couplings. This can be attributed to the fact that, in the latter case, the couplings are separately determined by the fits to the lepton and quark sector. As a result the individual c_i parameters in each sector are interlinked so as to reproduce the correct hierarchy in the mass matrix. For instance, given a choice c_{L_1} , there is less freedom in the choice of $c_{L_{2,3}}$. The λ'_{ijk} couplings on the other hand, depend on c_{L_i} , c_{Q_j} and c_{D_k} . Thus, for a given choice of c_{Q_j} , c_{D_k} , which are related from the quark mass fits, there is freedom in the choice of c_{L_i} which are determined from the fits to leptonic sector and are decoupled from the quark sector at tree level. We now proceed to discussing the phenomenological implications of the presence of these \mathcal{R}_p parameters.

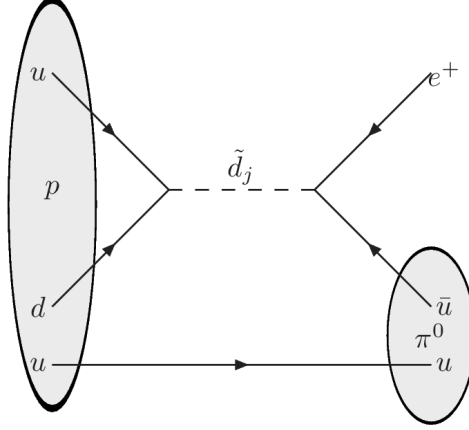


FIG. 3: Possible \mathcal{R}_p violating process ($p \rightarrow \pi^0 e^+$) yielding a non-zero decay rate for non-zero $\lambda'_{11k} \lambda''_{11j}$.

A. Nucleon decay

The presence of both lepton- and baryon-number violating terms in the lagrangian can give rise to small proton-decay lifetimes for baryon and lepton number violating couplings being simultaneously non-zero. For instance, a combination of $\lambda''_{ijk}, \lambda'_{ijk}$ can give rise to the contribution to proton decay shown in Fig. 3. This leads to particularly stringent constraints on the sizes of the couplings. Some of the strongest constraints come from searches for the following decays [43]:

$$\begin{aligned}
 |\lambda'_{l1k} \lambda''_{11k}| &\leq 2 \times 10^{-25} \left(\frac{\tilde{m}}{1 \text{ TeV}} \right)^2 \quad (l = 1, 2) \quad p \rightarrow [\pi^0 l^+] \\
 |\lambda'_{31k} \lambda''_{11k}| &\leq 7 \times 10^{-25} \left(\frac{\tilde{m}}{1 \text{ TeV}} \right)^2 \quad n \rightarrow [\pi^0 \bar{\nu}] \\
 |\lambda'_{i2k} \lambda''_{11k}| &\leq 3 \times 10^{-25} \left(\frac{\tilde{m}}{1 \text{ TeV}} \right)^2 \quad p \rightarrow [K^+ \bar{\nu}].
 \end{aligned} \tag{17}$$

There exist similar bounds on the product of lepton and baryon number violating couplings from other decay modes of the proton and neutron [44–47].

We note that, even with a relatively heavy sparticle spectrum at around 1 TeV, the product of the minimum values of the couplings in Fig. 2 will violate the bounds on nucleon decay. The violation of the bounds is more severe for the couplings involving third generation fermions. If one insists on simultaneous lepton- and baryon-number violation, then a choice of $\hat{\lambda}_{ijk}, \hat{\lambda}'_{ijk}, \hat{\lambda}''_{ijk} \sim \mathcal{O}(1)$ is no longer viable. Assuming one has a common scaling factor for each \mathcal{R}_p coupling, putting it equal to at most $\delta = 2 \times 10^{-4}$ is necessary to completely evade all the bounds from nucleon decay. Thus the nucleon decay problem is vastly ameliorated, but not solved, by warping. As a result, we do not pursue this further and instead focus on cases where either baryon-number *or*

lepton-number is violated, predicting a stable proton and none of the dangerous lepton and baryon number violating nucleon decay channels

B. Lepton Number Violation Only

Since baryon-number is conserved, the proton does not decay in this case. However in this scenario there are additional contributions to the neutrino mass: tree level contributions originating from μ_i [48], and loop-induced contributions from $\lambda_{ijk}, \lambda'_{ijk}$ [49, 50]. While it may be possible to generate $\mathcal{O}(0.1)$ eV neutrino masses with these couplings, it is very difficult to satisfy the solar and atmospheric neutrino data, which require neutrino mass splittings one or two orders of magnitude smaller than this. As a result, we focus on the parameter space where these contributions are suppressed in comparison to Dirac neutrino mass terms generated by Eq. 9. Non-supersymmetric Randall-Sundrum scenarios in which lepton-number violating effects could be hidden have been considered in Refs. [51, 52].

While it may be simple to suppress the masses of the electron and muon neutrinos, λ'_{333} (being relatively large compared to the other \mathcal{R}_p couplings) and give rise to too heavy a tau neutrino as compared to data. We focus on the following region of parameter space, which leads to \mathcal{R}_p contributions to neutrino masses that are not larger than the observed mass splittings:

$$\lambda'_{133} < 10^{-6}, \quad \lambda'_{233} < 5 \times 10^{-6}, \quad \lambda'_{333} < 5 \times 10^{-6}, \quad \mu_3 \lesssim 0.01 \text{ GeV}. \quad (18)$$

The condition Eq. 18 requires the lepton doublets to be far away from the IR brane. However in order to fit the required charged lepton masses, the SM singlets must then be localized relatively close to the IR brane. Picking one particular ‘good-fit’ point, we have:

$$c_{L_1} = 2.32, \quad c_{L_2} = 2.27, \quad c_{L_3} = 1.61, \quad c_{E_1} = 1.74, \quad c_{E_2} = 0.5, \quad c_{E_3} = 0.5. \quad (19)$$

The corresponding values of the lepton-number violating couplings in this case are represented by the points in Fig. 2. One may push the lepton doublets to be further away from the IR brane by choosing $c_{E_i} < 0.5$. However, this choice is not ideal as this may induce large off-diagonal elements in the slepton mass matrix, potentially leading to large (and excluded) flavour violation. Along with Eq. 19, one example of a good-fit point includes the following choices:

$$\begin{aligned} c_{Q_1} = 0.68, \quad c_{Q_2} = 1.04, \quad c_{Q_3} = 0.77, \quad c_{D_1} = 2.89, \quad c_{D_2} = 2.07, \quad c_{D_3} = 1.48, \\ c_{U_1} = 3.5, \quad c_{U_2} = 1.98, \quad c_{U_3} = 0.47, \quad M_1 = M_2 = 2.5 \text{ TeV}, \quad M_3 = 1.2 \text{ TeV}. \end{aligned} \quad (20)$$

With this choice, the masses of the neutrinos at the low scale can be determined using **SOFTSUSY** [53, 54] and are predicted to be:

$$m_{\nu_1} = 1.6 \times 10^{-6} \text{ eV}; \quad m_{\nu_2} = 7.4 \times 10^{-6} \text{ eV}; \quad m_{\nu_3} = 0.8 \text{ eV}. \quad (21)$$

These masses do respect the direct constraints upon neutrino masses, however they do not respect oscillation data, which require $\Delta m_{\text{atm}}^2 \sim 2 \times 10^{-3} \text{ eV}^2$ and $\Delta m_{\text{sol}}^2 \sim 7.5 \times 10^{-5} \text{ eV}^2$ to be the values of differences in the neutrino masses squared [55].

In order to suppress the \hat{R}_p contribution to the neutrino masses, we make the following choices:

$$\hat{\lambda}'_{133} = 0.1; \quad \hat{\lambda}'_{233} = 0.1; \quad \hat{\lambda}'_{333} = 0.2; \quad \hat{\mu}_3 = 0.1 \quad (22)$$

With this choice, the \hat{R}_p violating contributions to the neutrino masses are then

$$m_{\nu_1} = 1.0 \times 10^{-8} \text{ eV}, \quad m_{\nu_2} = 4.0 \times 10^{-6} \text{ eV}, \quad m_{\nu_3} = 0.008 \text{ eV}, \quad (23)$$

smaller than the values required to satisfy oscillation data. (We could also have suppressed the \hat{R}_p contribution by further raising the gaugino masses M_1, M_2 from the values in Eq. 20 at the expense of making the supersymmetric spectrum heavy, thus worsening the supersymmetric solution to the hierarchy problem).

In addition to the operators in Eq. 6, operators of the form $(L_i H_u)(L_j H_u)$ can also contribute to neutrino masses. This operator violates lepton number by 2 units. The superpotential term is given by

$$\mathcal{W}_{\Delta L=2} = \frac{\kappa_{ij}}{M_{Pl}} (L_i H_u) \cdot (L_j H_u), \quad (24)$$

where κ_{ij} are 5D Yukawa couplings with mass dimension M^{-1} . The neutrino mass matrix entry generated from this operator is given by

$$(m_\nu)_{ij} = \hat{\kappa}_{ij} \frac{v_u^2}{2M_{Pl}} e^{kR\pi} f(c_{L_i}) f(c_{L_j}) \quad (25)$$

where $\hat{\kappa}_{ij} = k\kappa_{ij}$ is a dimensionless $\mathcal{O}(1)$ parameter and the function f is defined in Eq. 10. For $c_{L_i} = c_{L_j} > 0.5$, this expression can be simplified to

$$(m_\nu)_{ij} \sim \frac{v_u^2}{2M_{Pl}} e^{(2-2c_{L_i})kR\pi}, \quad (26)$$

which for $c_{L_i} = c_{L_j} = 1.6$ comes out to be around 10^{-5} eV , much smaller than the Dirac mass contribution from Eq. 9. The $(L_i H_u)(L_j H_u)$ contribution is generally negligible in our model.

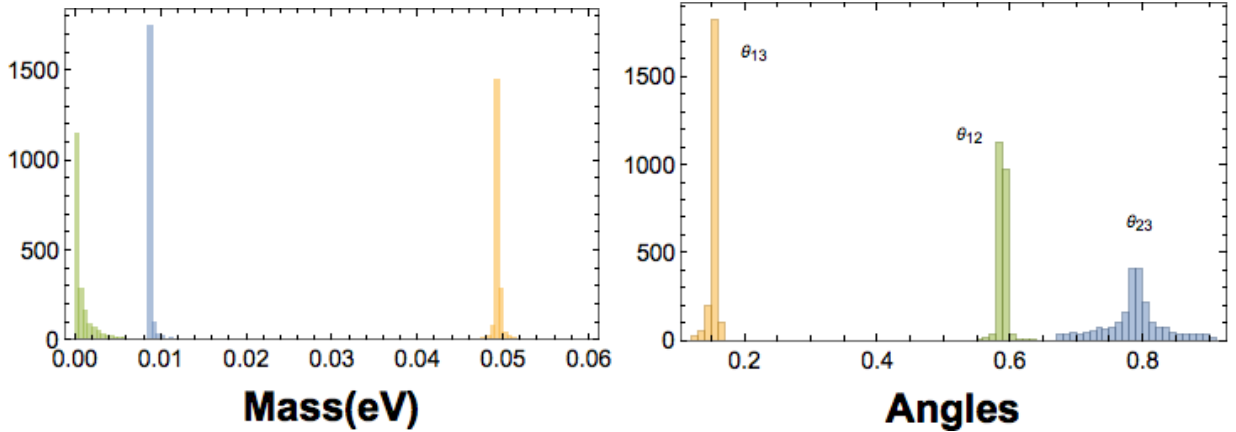


FIG. 4: The left hand plot shows the Dirac neutrino mass eigenvalues (eV) predicted by wave function overlap in RS models (a normal hierarchy is assumed). The right hand plot displays the predicted mixing angles. The vertical axis is effectively arbitrary, showing the frequency of the prediction in a large number of scanned models.

We are now free, after the addition of right-handed neutrino superfields, to arrange for dominant Dirac contributions to the neutrino masses. The oscillation parameters are determined from the fits to the leptonic data as outlined in Section II. The c_{N_i} parameters (for the right handed neutrinos) which pass the filtering criteria give rise to specific forms of neutrino mass textures leading to a determination of the neutrino parameters. The mixing angles and the mass eigenvalues can be determined by using the c_{N_i} in Eq.(9). Corresponding to the set in Eq.(19), the set of c_{N_i} parameters is:

$$c_{N_1} = 6.26 \quad c_{N_2} = 5.99 \quad c_{N_3} = 8.72 \quad (27)$$

Fig. 4 shows the results of the fit to the neutrino oscillation data obtained from the model parameters. The left-handed plot shows the predicted distribution of neutrino mass eigenvalues using different sets of c_{N_i} parameters satisfying both $\delta m_{12}^2 \sim 7.5 \times 10^{-5} \text{ eV}^2$ and $\delta m_{32}^2 \sim 2.32 \times 10^{-3} \text{ eV}^2$. In addition, the corresponding PMNS mixing angle predictions are shown in the right hand plot of Fig. 4, and they are close to the values inferred from experiment [55]. Thus we see that, predictions in line with oscillation data are easy to achieve in a RS model that generates the masses purely from wave function overlap in the Dirac masses.

The supersymmetric spectrum corresponding to the choice of GUT scale parameters⁶ obtained from Eqs. 19,20 is given in Table I. In our set-up, the higher order corrections to m_h are highly

⁶ It was shown in [56] that the running of soft masses may depend on physics in the hidden sector which breaks SUSY. These effects are likely to be relevant only for third generation squarks and are not included here because they add additional model dependence.

Parameter	Mass/TeV	Parameter	Mass/TeV	Parameter	Mass/TeV	Parameter	Mass/TeV
\tilde{t}_1	1.8	\tilde{b}_1	2.2	$\tilde{\tau}_1$	1.1	$\tilde{\nu}_\tau$	1.6
\tilde{t}_2	2.3	\tilde{b}_2	2.3	$\tilde{\tau}_2$	1.6	$\tilde{\nu}_\mu$	1.6
\tilde{c}_1	2.2	\tilde{s}_1	2.2	$\tilde{\mu}_R$	1.2	$\tilde{\nu}_e$	1.6
\tilde{c}_2	2.7	\tilde{s}_2	2.7	$\tilde{\mu}_L$	1.6	\tilde{g}	2.6
\tilde{u}_1	2.2	\tilde{d}_1	2.2	\tilde{e}_R	1.1	χ_1^\pm	2.0
\tilde{u}_2	2.7	\tilde{d}_2	2.7	\tilde{e}_L	1.6	χ_2^\pm	2.3
m_{A^0}	3.1	m_H^\pm	3.1	m_h	0.121	m_H	3.1
χ_1^0	1.1	χ_2^0	2.0	χ_3^0	2.3	χ_4^0	2.4

TABLE I: Example supersymmetric spectrum for the lepton number violating case and $\tan \beta = 5$.

suppressed as the masses of the KK excitations are 10^{16} GeV. Although the gravitational sector of the model is non-renormalizable, the higher order corrections are calculable to a good approximation because of this high suppression. SUSY kills any large corrections to the Higgs mass of order the Planck mass or GUT mass times a loop factor. The relevant higher order corrections are then from the renormalizable sector (the MSSM), which are calculable and of the order of a few GeV [57]. The lightest CP-even Higgs mass m_h is predicted a little on the low side, but the discrepancy with the experimental measurement of 0.125 TeV can be explained by missing higher order MSSM corrections in its prediction. The GUT-scale soft masses of the first two generations are typically of the order of 10-15 GeV at most, whereas the gaugino masses are of the order of 1 TeV, leading to running-induced approximate flavour degeneracy in the weak scale masses of the first two generations of sfermion.

The sparticle spectrum is heavy enough to have not yet been ruled out by LHC constraints, but light enough to expect a discovery in future runs. Indeed, the model predicts that there will be many multi-lepton rich signals from all of the lepton-number violating couplings that are switched on: strongly interacting sparticles are likely to be detected first. These then undergo cascade decay via R_p conserving processes until the lightest supersymmetric particle - in this case the \tilde{e}_R - is reached: this is because the R_p preserving dimensionless couplings such as gauge couplings and third family Yukawa couplings are larger than the \tilde{R}_p ones that are shown in Fig. 2. This particle

then decays via \tilde{R}_p : the predominant decay in this case is via λ'_{133} into a bottom quark and an anti-top. Thus, SUSY events are b -rich (predicting 4 b quarks) and may produce leptons from top decays, or be susceptible to top taggers. There is a non-zero branching ratio for $\tilde{e}_R \rightarrow \mu\nu_\tau$ via a similar sized coupling λ_{132} , and the additional muons may also aid detection strategies in multi-lepton channels.

We note that recent \tilde{R}_p explanations of an apparent excesses in LHC data [58–62] are not naturally accommodated in this set-up. They all are based on resonant slepton production and require a large order 0.1 coupling λ'_{i11} , which is not possible in our set-up. One would need additional flavour symmetry in order to fix some \tilde{R}_p couplings to zero and reduce the effect of various bounds on products of them in order to be able to accommodate such a coupling.

C. Baryon-number violation only

We now consider a scenario where only baryon-number violating terms are included in the lagrangian. Since lepton number is perturbatively conserved in this case, the superpotential terms proportional to λ_{ijk} , λ'_{ijk} or μ_i are absent. Proton decay is forbidden as it requires the presence of both baryon- and lepton-number violating terms in the Lagrangian. The neutrinos in this case must be purely Dirac type and their masses are determined using Eq. 9, just as for the other charged fermions. The results of the fit to the neutrino oscillation data is given in Fig. 4.

We illustrate the spectrum for the following parameter choice:

$$\begin{aligned} c_{Q_1} &= 2.2, \quad c_{Q_2} = 1.7, \quad c_{Q_3} = 0.7, \quad c_{D_1} = 1.8, \quad c_{D_2} = 1.2, \quad c_{D_3} = 1.4, \\ c_{U_1} &= 2.3, \quad c_{U_2} = 1.3, \quad c_{U_3} = 0.3, \quad c_{L_1} = 2.2, \quad c_{L_2} = 1.8, \quad c_{L_3} = 1.4, \\ c_{E_1} &= 1.7, \quad c_{E_2} = 0.9, \quad c_{E_3} = 0.5, \quad M_1 = 1.0 \text{ TeV}, \quad M_2 = 1.0 \text{ TeV}, \quad M_3 = 1.4 \text{ TeV}. \end{aligned} \tag{28}$$

The choice of the corresponding \tilde{R}_p couplings is represented by the blue points in Fig. 2. Note that the lepton doublet fields need not be so strongly localized towards the UV as in the lepton number violating case because there are no bilinear couplings which can contribute to the neutrino masses. This is reflected in the values of c_{L_i} . Table II gives the low energy spectrum corresponding to the choice of GUT scale parameters by Eq. 29. We find the spectrum has a nice feature wherein the coloured sparticles are grouped together in a small mass window. The sleptons in this case have a tendency to be lighter than the lepton number violating case as there are no constraints coming from upper bounds on the neutrino masses. The light smuon and neutralino gives a non-negligible contribution to the anomalous magnetic moment of the muon $(g - 2)_\mu$,

Parameter	Mass/TeV	Parameter	Mass/TeV	Parameter	Mass/TeV	Parameter	Mass/TeV
\tilde{t}_1	2.3	\tilde{b}_1	2.3	$\tilde{\tau}_1$	0.3	$\tilde{\nu}_\tau$	0.3
\tilde{t}_2	2.7	\tilde{b}_2	2.8	$\tilde{\tau}_2$	1.0	$\tilde{\nu}_\mu$	0.3
\tilde{c}_1	2.8	\tilde{s}_1	2.8	$\tilde{\mu}_1$	0.3	$\tilde{\nu}_e$	0.3
\tilde{c}_2	2.7	\tilde{s}_2	2.8	$\tilde{\mu}_2$	0.9	\tilde{g}	3.2
\tilde{u}_1	2.8	\tilde{d}_1	2.8	\tilde{e}_1	0.3	χ_1^\pm	0.8
\tilde{u}_2	2.6	\tilde{d}_2	2.8	\tilde{e}_2	0.9	χ_2^\pm	3.1
m_{A^0}	3.3	m_H^\pm	3.3	m_h	0.121	m_H	3.3
χ_1^0	0.1	χ_2^0	1.0	χ_3^0	2.0	χ_4^0	2.1

TABLE II: Example supersymmetric spectrum for $\tan \beta = 5$ in the baryon number violating case.

which may explain the apparent 3.6σ discrepancy between measurements and SM predictions: $\delta(g-2)_\mu/2 = (29 \pm 8) \times 10^{-10}$ [63]. SUSY loops with smuons and neutralinos running in the loop yield $\delta(g-2)_\mu/2 \approx 13 \times 10^{-10} \left(100 \text{ GeV}/\max(m_{\tilde{\mu}_L}, m_{\chi_1^0})\right)^2 \tan \beta$ [64]⁷. Thus, it appears that by increasing $\tan \beta$ (which may go as high as 50) one may fit $(g-2)_\mu/2$. Again, the spectrum presented is allowed by previous collider constraints, but should be covered in coming LHC runs. Again, production of the strongly interacting particles will proceed via R -parity conserving decays, and usually end in the lightest neutralino χ_1^0 . This will then decay via λ_{323}'' into a top, a strange and a bottom so we again expect bottom-rich events (at least four), but now there is no obvious source of missing energy unless leptons come from the top decay with an associated neutrino. The ‘golden’ decay chain $\tilde{q} \rightarrow \chi_2^0 q \rightarrow \tilde{e} e q \rightarrow \chi_1^0 e^+ e^- q$ is also open, which may lead to interesting invariant mass edges between the leptons (golden decays with e replaced by μ in the preceding decay should also be present).

IV. CONCLUSIONS

In a general supersymmetric extension of the SM, lepton and baryon number are not necessarily perturbatively conserved, unless a symmetry such as R -parity is invoked. As a result, the most

⁷ For recent work on explaining $g-2$ in \mathcal{R}_P SUSY see [65]

general supersymmetric lagrangian includes terms which violate both these symmetries. This however increases the number of free parameters in the form of undetermined values of the \mathcal{R}_p couplings.

In this work we propose a scenario by embedding the MSSM in a higher dimensional warped framework. Following the aesthetic that all dimensionless parameters should be of order 1 in a fundamental theory the warped dimensional set-up provides the flavour structure, while supersymmetry resolves the technical hierarchy problem. All of the supersymmetric parameters at the GUT scale including the \mathcal{R}_p couplings are determined by the same set of parameters which fix the fermion masses and mixings at this scale. This lends a certain level (order of magnitude-wise) of predictability to the framework, and we present, in Fig. 2, predictions of ranges of R -parity violating parameters. The couplings involving the third family tend to be the largest because of the warping structure. The predictions typically range over several orders of magnitude but are dependent on the flavour indices of the coupling.

In fact, in Ref. [34], a related approach was taken by examining R -parity violation in a *TeV-scale* RS model (as opposed to our *GUT-scale* RS set-up). In the TeV-scale RS case, bounds from proton decay are much more stringent because a dimension-5 operator $QQQL$ can be generated on the IR brane, suppressed only by the TeV scale. This was surmounted by imposing a lepton-number triality. The resulting baryon-number violating couplings involving the up and charm quarks were around 0.02-0.05, whereas those involving tops were highly suppressed, primarily because the couplings were dominantly generated by the RPV operators localized on the UV brane. The predictions of the TeV-scale RS case are in contrast to ours in Fig. 2, which shows that the couplings involving tops are the largest, but these may be as large as $\mathcal{O}(10^{-4}) - \mathcal{O}(10^{-2})$. The magnitude of these couplings can be attributed to the fact that in our set-up, the Higgs fields, SUSY breaking spurion as well as the RPV operators are localized on the same brane. This was motivated from the requirement of having a model consistent with FCNC constraints in addition to having light sfermions.

Returning now to our GUT-scale RS scenario, for the most general case which includes both baryon and lepton number violating terms, the nucleon decays too quickly for $\sim \mathcal{O}(1)$ dimensionless \mathcal{R}_p parameters, although if instead they are all set to be $\mathcal{O}(10^{-4})$, the lifetime may be long enough to evade current experimental bounds (for superpartner masses of around 2 TeV). Following our initial idea of the aesthetic, it appears though that one needs to forbid either the lepton-number or baryon-number violating terms, in which case plenty of parameter space exists where current experimental bounds on the couplings are respected.

For the case where only lepton number is violated we find points in parameter space where the neutrino masses are predominantly Dirac-like nature, even in the presence of various lepton number violating operators contributing to the neutrino masses. The neutrino masses and mixings are fit to oscillation data just as the charged fermions are fit. In the baryon number violating case, the sleptons have a tendency to be lighter making it more appealing from the collider searches point of view: leptons may appear more often in supersymmetric decay chains, providing clean objects with low backgrounds to search for. In addition, the lighter smuons mean that a supersymmetric explanation for the discrepant anomalous magnetic moment of the muon is viable. In either the lepton-number or baryon-number violating cases, LHC signals consist of prompt hard jets, and b -rich events (at least four per event are predicted) containing tops. In the lepton-number violating case there may also be a modest amount of missing transverse momentum coming from neutrino production. We illustrate points in parameter space where current collider limits are respected but where the LHC should be able to discover sparticles in future runs, which we eagerly await.

V. ACKNOWLEDGEMENTS

This work was partially supported by STFC grant ST/L000385/1. BCA was adjunct faculty in DTP TIFR during December 2014 and would like to thank the department for the hospitality extended during the early stages of this work and the Cambridge SUSY Working Group for helpful suggestions. BCA, AI and KS would like to thank Sudhir Vempati for the motivation and useful discussions. AI and KS also thank the organisers at WHEPP 2013 for the hospitality where the idea was conceived. AI and KS would also like to thank the CERN theory division for hospitality where part of the work was completed. AI would like to thank Tuhin Roy for discussions on contribution to renormalization of soft masses due to hidden sector effects.

AI would like to dedicate this paper to the memory of close friend and fellow physicist Dr. Senti Imsong, who passed away recently.

-
- [1] S. Chatrchyan et al. (CMS), Phys. Lett. **B716**, 30 (2012), 1207.7235.
 - [2] G. Aad et al. (ATLAS), Phys. Lett. **B716**, 1 (2012), 1207.7214.
 - [3] S. Chatrchyan et al. (CMS), JHEP **06**, 081 (2013), 1303.4571.
 - [4] A. E. Nelson and M. J. Strassler, JHEP **09**, 030 (2000), hep-ph/0006251.

- [5] E. Dudas, G. von Gersdorff, J. Parmentier, and S. Pokorski, JHEP **12**, 015 (2010), 1007.5208.
- [6] C. Froggatt and H. B. Nielsen, Nucl.Phys. **B147**, 277 (1979).
- [7] G. R. Farrar and P. Fayet, Phys. Lett. **B76**, 575 (1978).
- [8] S. Dimopoulos and H. Georgi, Nucl. Phys. **B193**, 150 (1981).
- [9] N. Sakai and T. Yanagida, Nucl. Phys. **B197**, 533 (1982).
- [10] S. Dimopoulos, S. Raby, and F. Wilczek, Phys. Lett. **B112**, 133 (1982).
- [11] S. Weinberg, Phys. Rev. **D26**, 287 (1982).
- [12] L. M. Krauss and F. Wilczek, Phys. Rev. Lett. **62**, 1221 (1989).
- [13] A. Font, L. E. Ibanez, and F. Quevedo, Phys. Lett. **B228**, 79 (1989).
- [14] R. N. Mohapatra, Phys. Rev. **D34**, 909 (1986).
- [15] S. P. Martin, Phys. Rev. **D46**, 2769 (1992), hep-ph/9207218.
- [16] S. P. Martin, Phys. Rev. **D54**, 2340 (1996), hep-ph/9602349.
- [17] B. C. Allanach, A. Dedes, and H. K. Dreiner, Phys. Rev. **D69**, 115002 (2004), [Erratum: Phys. Rev.D72,079902(2005)], hep-ph/0309196.
- [18] R. Barbier et al., Phys. Rept. **420**, 1 (2005), hep-ph/0406039.
- [19] F. Takayama and M. Yamaguchi, Phys. Lett. **B485**, 388 (2000), hep-ph/0005214.
- [20] G. Moreau and M. Chemtob, Phys. Rev. **D65**, 024033 (2002), hep-ph/0107286.
- [21] T. Gherghetta and A. Pomarol, Nucl. Phys. **B586**, 141 (2000), hep-ph/0003129.
- [22] R. Contino and A. Pomarol, JHEP **11**, 058 (2004), hep-th/0406257.
- [23] T. Gherghetta, in *Physics of the large and the small, TASI 09, proceedings of the Theoretical Advanced Study Institute in Elementary Particle Physics, Boulder, Colorado, USA, 1-26 June 2009* (2011), pp. 165–232, 1008.2570, URL <https://inspirehep.net/record/865392/files/arXiv:1008.2570.pdf>.
- [24] L. Randall and R. Sundrum, Phys. Rev. Lett. **83**, 3370 (1999), hep-ph/9905221.
- [25] D. Marti and A. Pomarol, Phys. Rev. **D64**, 105025 (2001), hep-th/0106256.
- [26] K.-w. Choi, D. Y. Kim, I.-W. Kim, and T. Kobayashi, Eur. Phys. J. **C35**, 267 (2004), hep-ph/0305024.
- [27] K.-w. Choi, D. Y. Kim, I.-W. Kim, and T. Kobayashi (2003), hep-ph/0301131.
- [28] F. Brummer, S. Fichet, and S. Kraml, JHEP **12**, 061 (2011), 1109.1226.
- [29] A. M. Iyer and S. K. Vempati, Phys. Rev. **D88**, 016005 (2013), 1304.3558.
- [30] S. Chang, J. Hisano, H. Nakano, N. Okada, and M. Yamaguchi, Phys. Rev. **D62**, 084025 (2000), hep-ph/9912498.
- [31] Y. Grossman and M. Neubert, Phys. Lett. **B474**, 361 (2000), hep-ph/9912408.
- [32] A. M. Iyer and S. K. Vempati, Phys. Rev. **D86**, 056005 (2012), 1206.4383.
- [33] W. D. Goldberger, Y. Nomura, and D. Tucker-Smith, Phys. Rev. **D67**, 075021 (2003), hep-ph/0209158.
- [34] B. Heidenreich and Y. Nakai, JHEP **10**, 182 (2014), 1407.5095.
- [35] S. Raychaudhuri and K. Sridhar, *Particle Physics of Brane Worlds and Extra Dimensions* (Cambridge University Press, to appear).
- [36] F. James and M. Roos, Comput. Phys. Commun. **10**, 343 (1975).

- [37] A. de Gouvea, S. Lola, and K. Tobe, Phys. Rev. **D63**, 035004 (2001), hep-ph/0008085.
- [38] D. Choudhury and P. Roy, Phys. Lett. **B378**, 153 (1996), hep-ph/9603363.
- [39] J. L. Goity and M. Sher, Phys. Lett. **B346**, 69 (1995), [Erratum: Phys. Lett.B385,500(1996)], hep-ph/9412208.
- [40] G. Bhattacharyya, D. Choudhury, and K. Sridhar, Phys. Lett. **B349**, 118 (1995), hep-ph/9412259.
- [41] G. Bhattacharyya, J. R. Ellis, and K. Sridhar, Mod. Phys. Lett. **A10**, 1583 (1995), hep-ph/9503264.
- [42] F. Ledroit and G. Sajot, Rapport GDR-Supersymetrie, GDR-S-008, ISN, Grenoble. (1998).
- [43] S. Eidelman et al. (Particle Data Group), Phys. Lett. **B592**, 1 (2004).
- [44] K. Rajagopal, M. S. Turner, and F. Wilczek, Nucl. Phys. **B358**, 447 (1991).
- [45] D. Chang and W.-Y. Keung, Phys. Lett. **B389**, 294 (1996), hep-ph/9608313.
- [46] G. Bhattacharyya and P. B. Pal, Phys. Lett. **B439**, 81 (1998), hep-ph/9806214.
- [47] G. Bhattacharyya and P. B. Pal, Phys. Rev. **D59**, 097701 (1999), hep-ph/9809493.
- [48] A. S. Joshipura and M. Nowakowski, Phys. Rev. **D51**, 2421 (1995), hep-ph/9408224.
- [49] L. J. Hall and M. Suzuki, Nucl. Phys. **B231**, 419 (1984).
- [50] K. S. Babu and R. N. Mohapatra, Phys. Rev. Lett. **64**, 2747 (1990).
- [51] T. Gherghetta, Phys. Rev. Lett. **92**, 161601 (2004), hep-ph/0312392.
- [52] A. M. Iyer and S. K. Vempati, Phys. Rev. **D88**, 073005 (2013), 1307.5773.
- [53] B. C. Allanach, Comput. Phys. Commun. **143**, 305 (2002), hep-ph/0104145.
- [54] B. C. Allanach, C. H. Kom, and M. Hanussek, Comput. Phys. Commun. **183**, 785 (2012), 1109.3735.
- [55] M. Maltoni, T. Schwetz, M. A. Tortola, and J. W. F. Valle, New J. Phys. **6**, 122 (2004), hep-ph/0405172.
- [56] A. G. Cohen, T. S. Roy, and M. Schmaltz, JHEP **02**, 027 (2007), hep-ph/0612100.
- [57] T. Hahn, S. Heinemeyer, W. Hollik, H. Rzehak, and G. Weiglein, PoS **LL2014**, 032 (2014), 1407.0232.
- [58] B. Allanach, S. Biswas, S. Mondal, and M. Mitra, Phys. Rev. **D91**, 011702 (2015), 1408.5439.
- [59] B. C. Allanach, S. Biswas, S. Mondal, and M. Mitra, Phys. Rev. **D91**, 015011 (2015), 1410.5947.
- [60] B. C. Allanach, P. S. B. Dev, and K. Sakurai (2015), 1511.01483.
- [61] R. Ding, L. Huang, T. Li, and B. Zhu (2015), 1512.06560.
- [62] B. C. Allanach, P. S. B. Dev, S. A. Renner, and K. Sakurai (2015), 1512.07645.
- [63] G. W. Bennett et al. (Muon g-2), Phys. Rev. **D73**, 072003 (2006), hep-ex/0602035.
- [64] A. Czarnecki and W. J. Marciano, Phys. Rev. **D64**, 013014 (2001), hep-ph/0102122.
- [65] A. Chakraborty and S. Chakraborty (2015), 1511.08874.